

Photoproduction of hypernuclei within the quark-meson coupling model

R. Shyam^{a,c}, K. Tsushima^b, A. W. Thomas^a

^a*Theory Center, Thomas Jefferson National Accelerator Facility, Newport News, VA 23606, USA*

^b*Excited Baryon Analysis Center and Theory Center, Thomas Jefferson National Accelerator Facility, Newport News, VA 23606, USA*

^c*Theory Division, Saha Institute of Nuclear Physics, Kolkata 700064, India*

Abstract

We study the photoproduction of the $^{12}\Lambda\text{B}$ hypernucleus within a fully covariant effective Lagrangian based model, employing Λ bound state spinors derived from the latest quark-meson coupling model. The kaon production vertex is described via creation, propagation and decay of $N^*(1650)$, $N^*(1710)$, and $N^*(1720)$ intermediate baryonic resonant states in the initial collision of the photon with a target proton in the incident channel. The parameters of the resonance vertices are fixed by describing the total and differential cross section data on the elementary $\gamma p \rightarrow \Lambda K^+$ reaction in the energy regime relevant to the hypernuclear production. It is found that the hypernuclear production cross sections calculated with the quark model based hyperon bound state spinors differ significantly from those obtained with the phenomenological Dirac single particle wave functions.

Key words: Photoproduction of hypernuclei, covariant production model, quark-meson coupling model hyperon spinors.

PACS: 21.80.+a, 13.60.-r, 13.75.Jz

Electromagnetic probes provide a very powerful tool for studying the Λ hypernuclei. In contrast to the hadronic reactions [(K^-, π^-) and (π^+, K^+)], a proton in the target nucleus is converted into a Λ hyperon in both (γ, K^+) and $(e, e' K^+)$ reactions, thus forming a neutron-rich hypernucleus. This leads to the formation of mirror hypernuclear systems which can facilitate the study of the charge symmetry breaking with strangeness degrees of freedom (see, e.g., [1, 2, 3, 4]). Although in the electromagnetic reactions the momentum

transfer to the nucleus is comparable to that of the (π^+, K^+) reaction, they carry, in addition, significant spin-flip amplitudes due to the absorption of the photon spin and the forward angle domination of the cross sections. Furthermore, while the hadronic hypernuclear production reactions are confined mostly to the nuclear surface because of strong absorption of both K^- and π^\pm , the electromagnetic reactions occur deep inside the nucleus because of the weaker nuclear interactions of both photon and K^+ . This makes them an ideal tool for studying deeply bound hypernuclear states.

Recently, Jefferson Laboratory (JLab) has started a systematic study of the high-resolution hypernuclear production reactions on p -shell target nuclei (${}^9\text{Be}$, ${}^{12}\text{C}$ and ${}^{16}\text{O}$) using continuous electron beams [5, 6, 7, 8]. The quality and high resolution (~ 400 keV) of the electron beam in these experiments make it possible to identify hyperon single particle states more clearly and to untangle the core excited states for the first time. While, the first measurements of hypernuclear production with real photons [(γ, K^+) reaction] on a nuclear target (${}^{12}\text{C}$) were reported long ago [9], interest in this field has been revived with the possibility of performing more such measurements at accelerators MAMI-C in Mainz, and ELSA in BONN (see, e.g. [10]).

Several theoretical studies of photoproduction of hypernuclei have been reported [12, 13, 14, 15, 16, 17]. They all use the framework of the impulse approximation, where the hypernuclear production amplitudes are calculated by determining expectation values of the operator for the elementary $p(\gamma, K^+)\Lambda$ process. This operator is constructed either by using the Feynman diagrammatic approach including graphs corresponding to Born terms and resonance terms in s and u channels [12, 13, 17, 18], or phenomenologically by parameterizing the experimental cross sections for the elementary process [15, 16]. Except for Ref. [14], where Dirac spinors were used to describe the initial and final bound state wave functions, nonrelativistic models have been employed to obtain these wave functions in all of these investigations.

On the other hand, in Ref. [19] a fully covariant model was employed to calculate the cross sections of the ${}^{16}\text{O}(\gamma, K^+){}^{16}_{\Lambda}\text{N}$ reaction. This model retains the full field theoretical structure of the interaction vertices and treats the baryons as Dirac particles (see also Ref. [20]). The initial state interaction of the incoming photon with a bound proton leads to excitations of $N^*(1650)$ [$\frac{1}{2}^-$], $N^*(1710)$ [$\frac{1}{2}^+$], and $N^*(1720)$ [$\frac{3}{2}^+$] resonance intermediate states, which have been shown to make the predominant contributions to the $p(\gamma, K^+)\Lambda$ cross section [11]. In this model calculations are performed in momentum

space throughout, hence it includes all the nonlocalities in the production amplitudes that arise from the resonance propagators.

However, the procedure of obtaining the bound state spinors in the previous application of this model brings in some uncertainty in the calculated hypernuclear production cross sections. In Ref. [19] the bound state spinors were computed in the coordinate space by solving the Dirac equation with scalar and vector fields having a Woods-Saxon radial form. With a set of radius and diffuseness parameters, the depths of these fields are searched so as to reproduce the binding energy (BE) of the given state. Because the experimental BEs of the hypernuclear states often involve ambiguities, the extracted potential depths also become ambiguous. Besides, the depths of the potential fields are dependent on the adopted radius and diffuseness parameters and there is no certain way of fixing them. Furthermore, both vector and scalar fields are assumed to have the same geometry.

In this paper, we explore the feasibility of studying the photoproduction of hypernuclei within the relativistic model of Ref. [19] but employing hyperon bound state spinors calculated within the quark-meson coupling (QMC) model. This provides an opportunity to investigate the role of the quark degrees of freedom in the hypernuclear production, which is a novel feature of this study. Since photoproduction of hypernuclei involves large momentum transfers [21] to the target nucleus, it appears to be a good case for examining such short distance effects.

In the QMC model [22, 23, 24, 25], quarks in the non-overlapping bags (modeled using MIT bag), interact self consistently with isoscalar-scalar (σ) and isoscalar-vector (ω) mesons in the mean field approximation. The explicit treatment of the nucleon internal structure represents an important departure from quantum hadro-dynamics (QHD) model [26]. The self-consistent response of the bound quarks to the mean σ field leads to a new saturation mechanism for nuclear matter [22]. The QMC model has been used to study the properties of finite nuclei [27], the binding of ω , η , η' and D nuclei [28, 29, 30, 31] and also the effect of the medium on K^\pm and J/Ψ production [32].

The most recent development of the quark-meson coupling model is the inclusion of the self-consistent effect of the mean scalar field on the familiar one-gluon exchange hyperfine interaction that in free space leads to the $N - \Delta$ and $\Sigma - \Lambda$ mass splitting [33]. With this [34] the QMC model has been able to explain the properties of Λ hypernuclei for the s -states rather well, while the p - and d -states tend to underbind. It also leads to a very natural

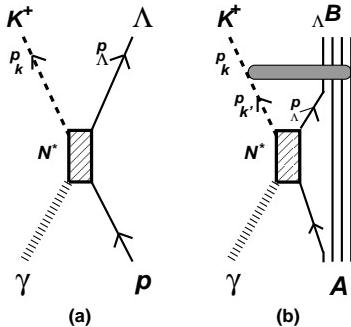


Figure 1: Representation of the type of Feynman diagrams included in our calculations. The elliptic shaded area represents the optical model interactions in the outgoing channel.

explanation of the small spin-orbit force in Λ -nucleus interaction. In this exploratory work, the bound Λ spinors are generated from this version of the QMC model and are used to calculate the cross sections of the $^{12}\text{C}(\gamma, K^+)^{12}\text{B}$ reaction.

We fix the parameters of the resonance vertices by describing the data on total and differential cross sections of the elementary $\gamma p \rightarrow K^+\Lambda$ process in the relevant photon energy regime within a similar effective Lagrangian approach. This is in contrast to the calculations presented in Ref. [19] where they were taken from previous studies of photon and hadron induced associated $K^+\Lambda$ production reactions [11, 35, 36]. Thus the resonance parameters used in the present study are better constrained.

A preliminary experimental investigation of the (γ, K^+) reaction on ^{12}C was reported already in 1995 [9]. Recently, the ^{12}B hypernucleus has been produced at JLab via the $(e, e'K^+)$ reaction with a very high energy resolution [5, 7]. In this experiment, apart from observing hypernuclear excitations where a proton is replaced (leaving ^{11}B in the $3/2^-$ ground state) by a Λ in s and p shells, one also sees identifiable strength in the region which corresponds to the excitation of the ^{11}B core. This underlines the need of using a more microscopic hypernuclear structure model in describing the excitation of hypernuclear spectra in electromagnetic reactions. Our work is a first step in this direction where we examine the differences between the hypernuclear photoproduction cross sections obtained with a microscopic hypernuclear structure model and a phenomenological model. We restrict ourselves to photon energies below 1.5 GeV as this is the relevant energy regime for the experiment performed already with real photons [9]. Moreover, it has

been shown in previous studies [14, 16, 19] that hypernuclear photoproduction cross sections on light targets peak around 0.95 - 1.0 GeV and drop off thereafter.

As in the previous study [19], we have used the graph of the type shown in Fig. 1(b) to describe the hypernuclear production reaction $A(\gamma, K^+)_{\Lambda}B$. The elementary $\gamma p \rightarrow K^+ \Lambda$ process involved in this diagram is shown in Fig. 1(a). It is clear that our model has only *s*-channel resonance contributions. In principle, Born terms and resonance contributions in *u*- and *t*-channels should also be included in the description of both the processes. These graphs constitute the non-resonant background contributions. It should be noted that their magnitudes depend on the particular model used to calculate them and also on the parameters used within that model [37]. Except for photon energies close to threshold, these terms have been shown to make non-negligible contributions in the models of Refs. [11, 37]. On the other hand, in Ref. [17] they have been found to be insignificant in both elementary as well as in-medium photon induced reactions for beam energies below 1.5 GeV. We have ignored these diagrams in this exploratory work to keep our production model simple and similar to that of Ref. [19]. Furthermore, we reduce the computational complications further by using plane waves (PW) to describe the relative motion of the outgoing particle which is justified by the relatively weaker kaon-nucleus interaction in the final channel.

All the ingredients (effective Lagrangians, resonance propagators etc.) required in calculations of the amplitudes associated with diagrams 1(a) and 1(b) are described in Refs. [19, 20]. The coupling constants for the $N^*N\gamma$ ($g_{p\gamma}^{1,2}$) and $N^*K^+\Lambda$ ($g_{K\Lambda}$) vertices used in the present study (shown in table I) have been determined by comparing our calculations [graph 1(a)] with the total and differential cross section data for the elementary $\gamma p \rightarrow \Lambda K^+$ reaction in the relevant photon energy region. In Figs. 2(a) and 2(b) we have shown the comparison of our results with the experimental data. While in Fig. 2(b) we have shown the data from both SAPHIR [39] and CLAS [40] collaborations, only SAPHIR data are shown in Fig. 2(a) as the CLAS total cross sections are somewhat uncertain due to absence of complete angular coverage. As was noted before [40], in the energy range of our interest the CLAS and SAPHIR data agree with each other fairly well. Nevertheless, the CLAS data have much less statistical uncertainties. We see that calculated cross sections are in close agreement with the data in the considered photon energy regime.

We further note that within our model $N^*(1650)$ resonance makes the

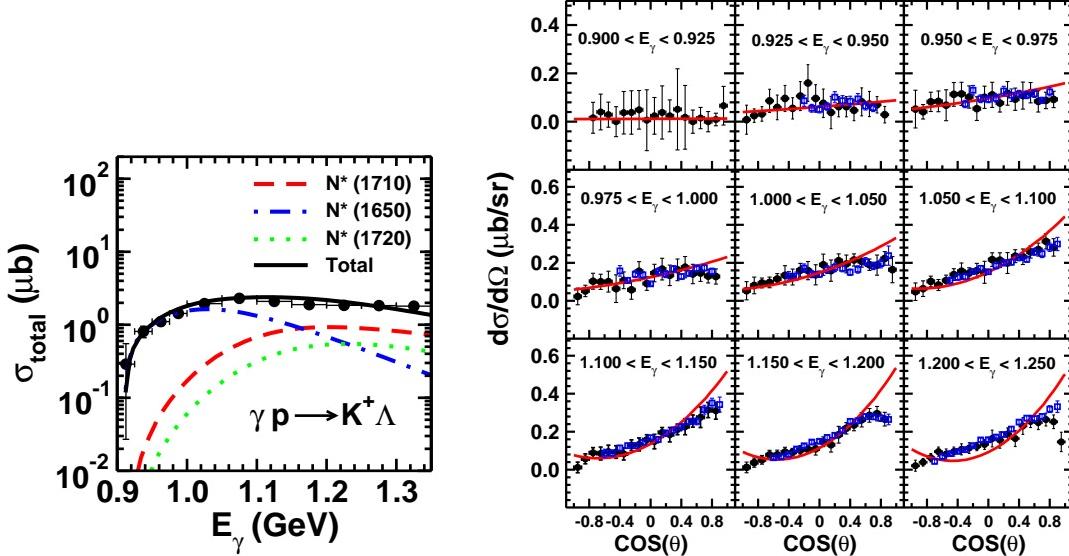


Figure 2: Calculated total cross sections (left panel) and differential cross section (right panel) for the $\gamma p \rightarrow \Lambda K^+$ reaction, obtained with the vertex constants of Table I as compared with the data from [39] (solid circles) and [40] (open squares).

dominant contribution to the total cross section at lower photon energies, while $N^*(1710)$ is more important at higher energies. The contribution of $N^*(1720)$ is much weaker everywhere. We add however that this result is particular to our single channel model. In calculations where more resonances and the background terms are included, the pattern of relative resonance contributions could be different due to e.g. different kind of interference effects. Indeed the $N^*(1710)$ contribution has been found to be weak in several of the recent models [37, 41, 42]. In the unitary coupled channels calculations of Ref. [11] this resonance is suppressed because of the destructive interference with the background terms. However, since our purpose was to fix the parameters at the hyperon production vertices in the hypernuclear production reaction considered in Fig. (2b), we consider the diagram 1(a) for the elementary reaction to be adequate.

The amplitudes for diagram 1(b) involve momentum space four component (spin space) Dirac spinors [$\psi(p)$] for bound nucleon and hyperon states [43] and the momentum space kaon-nucleus wave function [$\phi_K^{(-)*}(p'_K, p_K)$] which can be calculated by using an appropriate K^+ - nucleus optical potential (see, e.g., Ref. [44]). Momenta p_K and p'_K are as defined in Fig. 1(b). In

Table 1: Resonances included in the calculations and their coupling constants

Resonance	mass GeV	width GeV	$g_{p\gamma}^1$	$g_{p\gamma}^2$	$g_{K\Lambda}$
$N^*(1650)$	1.650	0.165	-0.45		0.96
$N^*(1710)$	1.710	0.180	0.25		-6.10
$N^*(1720)$	1.720	0.200	-0.75	0.25	0.07

the PW approximation one writes $\Phi_K^{(-)*}(p'_K, p_K) = \delta^4(p'_K - p_K)$.

The spinors, $\psi(p)$, are solutions of the Dirac equation in momentum space for a bound state problem in the presence of an external potential field [20, 43]

$$p\psi(p) = m_N\psi(p) + F(p), \quad (1)$$

where

$$\begin{aligned} F(p) &= \delta(p_0 - E) \left[\int d^3 p' V_s(-\mathbf{p}') \psi(\mathbf{p} + \mathbf{p}') \right. \\ &\quad \left. - \gamma_0 \int d^3 p' V_v^0(-\mathbf{p}') \psi(\mathbf{p} + \mathbf{p}') \right]. \end{aligned} \quad (2)$$

In Eq. (2), the real scalar and timelike vector potentials V_s and V_v^0 represent, respectively, the momentum space local Lorentz covariant interaction of single nucleon or Λ with the remaining $(A - 1)$ nucleons. We denote a four momentum by $p = (p_0, \mathbf{p})$. The magnitude of the three momentum \mathbf{p} is represented by k , and its directions by \hat{p} . p_0 is the time like component of p . Spinors $\psi(p)$ and $F(p)$ are written as

$$\begin{aligned} \psi(p) &= \delta(p_0 - E) \begin{pmatrix} f(k) \mathcal{Y}_{\ell 1/2j}^{m_j}(\hat{p}) \\ -ig(k) \mathcal{Y}_{\ell' 1/2j}^{m_j}(\hat{p}) \end{pmatrix}, \\ F(p) &= \delta(p_0 - E) \begin{pmatrix} \zeta(k) \mathcal{Y}_{\ell 1/2j}^{m_j}(\hat{p}) \\ -i\zeta'(k) \mathcal{Y}_{\ell' 1/2j}^{m_j}(\hat{p}) \end{pmatrix}, \end{aligned} \quad (3)$$

where $f(k)[\zeta(k)]$ is the radial part of the upper component of the spinor $\psi(p)[F(p)]$. Similarly $g(k)[\zeta'(k)]$ are the same of their lower component. $f(k)$ and $g(k)$ represent Fourier transforms of radial parts of the corresponding

Table 2: Depths of the Dirac vector (V_v) and scalar (V_s) fields for single particle Λ and nucleon shells. In each case, both fields have the Woods-Saxon form with similar radius ($r = 0.983$ fm) and diffuseness ($a = 0.606$ fm) parameters. Also shown are the experimental binding energies for each shell (numbers in the brackets are the BEs predicted by the QMC model).

State	BE (MeV)	V_v (MeV)	V_s (MeV)
$^{12}_\Lambda B(1s_{1/2})$	11.37 (14.93)	171.78	-212.69
$^{12}_\Lambda B(1p_{3/2})$	1.37 (3.62)	204.16	-252.28
$^{12}_\Lambda B(1p_{1/2})$	1.03 (3.62)	227.83	-280.86
$^{12}C(1p_{3/2})$	15.96	382.60	-472.34

coordinate space spinors. $\zeta(k)$ are related to f , g and the scalar and vector potentials as shown in Ref. [20].

In Table 2 we show the parameters associated with the scalar and vector fields of the phenomenological model for Λ and nucleon bound states, and the corresponding experimental BEs which are the averages of the values reported by several experimental studies [5, 6, 7, 45]. In this table we also give the BEs of the Λ bound states as predicted by the QMC model.

To calculate the bound state spinors within the QMC model we have used its latest version, where the calculations for Λ and Ξ hypernuclei are of comparable quality to earlier QMC results [29]. In addition, without requiring any additional parameter it predicts no nuclear bound Σ states [34], which is in qualitative agreement with the experimental observations. This is facilitated by the extra repulsion associated with the increased one-gluon-exchange hyperfine in-medium interaction. We refer to Ref. [34] for more details of this new version of the QMC.

In order to calculate the properties of finite hypernuclei, we construct a simple, relativistic shell model, with the nucleon core calculated in a combination of self-consistent scalar and vector mean fields. The Lagrangian density for a hypernuclear system in the QMC model is written as a sum of

two terms, $\mathcal{L}_{QMC}^{HY} = \mathcal{L}_{QMC} + \mathcal{L}_{QMC}^Y$, where [28],

$$\begin{aligned}\mathcal{L}_{QMC} = & \bar{\psi}_N(\vec{r})[i\gamma \cdot \partial - M_N(\sigma) - (g_\omega \omega(\vec{r}) \\ & + g_\rho \frac{\tau_3^N}{2} b(\vec{r}) + \frac{e}{2}(1 + \tau_3^N) A(\vec{r})) \gamma_0] \psi_N(\vec{r}) \\ & - \frac{1}{2}[(\nabla \sigma(\vec{r}))^2 + m_\sigma^2 \sigma(\vec{r})^2] \\ & + \frac{1}{2}[(\nabla \omega(\vec{r}))^2 + m_\omega^2 \omega(\vec{r})^2] \\ & + \frac{1}{2}[(\nabla b(\vec{r}))^2 + m_\rho^2 b(\vec{r})^2] + \frac{1}{2}(\nabla A(\vec{r}))^2,\end{aligned}\quad (4)$$

$$\begin{aligned}\mathcal{L}_{QMC}^Y = & \sum_{Y=\Lambda, \Sigma, \Xi} \bar{\psi}_Y(\vec{r})[i\gamma \cdot \partial - M_Y(\sigma) - (g_\omega^Y \omega(\vec{r}) \\ & + g_\rho^Y I_3^Y b(\vec{r}) + e Q_Y A(\vec{r})) \gamma_0] \psi_Y(\vec{r}),\end{aligned}\quad (5)$$

where $\psi_N(\vec{r})$ ($\psi_Y(\vec{r})$) and $b(\vec{r})$ are, respectively, the nucleon (hyperon) and the ρ meson (the time component in the third direction of isospin) fields, while m_σ , m_ω and m_ρ are the masses of the σ , ω and ρ mesons. g_ω and g_ρ are the ω -N and ρ -N coupling constants which are related to the corresponding (u,d)-quark- ω , g_ω^q , and (u, d) quark- ρ , g_ρ^q , coupling constants as $g_\omega = 3g_\omega^q$ and $g_\rho = g_\rho^q$. I_3^Y and Q_Y are the third component of the hyperon isospin operator and its electric charge in units of the proton charge, e , respectively.

The following set of equations of motion are obtained for the hypernuclear system from the Lagrangian density Eqs. (4)-(5):

$$\begin{aligned}[i\gamma \cdot \partial - M_N(\sigma) - (g_\omega \omega(\vec{r}) + g_\rho \frac{\tau_3^N}{2} b(\vec{r}) \\ + \frac{e}{2}(1 + \tau_3^N) A(\vec{r})) \gamma_0] \psi_N(\vec{r}) = 0,\end{aligned}\quad (6)$$

$$\begin{aligned}[i\gamma \cdot \partial - M_Y(\sigma) - (g_\omega^Y \omega(\vec{r}) + g_\rho^Y I_3^Y b(\vec{r}) \\ + e Q_Y A(\vec{r})) \gamma_0] \psi_Y(\vec{r}) = 0,\end{aligned}\quad (7)$$

$$(-\nabla_r^2 + m_\sigma^2) \sigma(\vec{r}) = g_\sigma C_N(\sigma) \rho_s(\vec{r}) + g_\sigma^Y C_Y(\sigma) \rho_s^Y(\vec{r}), \quad (8)$$

$$(-\nabla_r^2 + m_\omega^2) \omega(\vec{r}) = g_\omega \rho_B(\vec{r}) + g_\omega^Y \rho_B^Y(\vec{r}), \quad (9)$$

$$(-\nabla_r^2 + m_\rho^2) b(\vec{r}) = \frac{g_\rho}{2} \rho_3(\vec{r}) + g_\rho^Y I_3^Y \rho_B^Y(\vec{r}), \quad (10)$$

$$(-\nabla_r^2) A(\vec{r}) = e \rho_p(\vec{r}) + e Q_Y \rho_B^Y(\vec{r}), \quad (11)$$

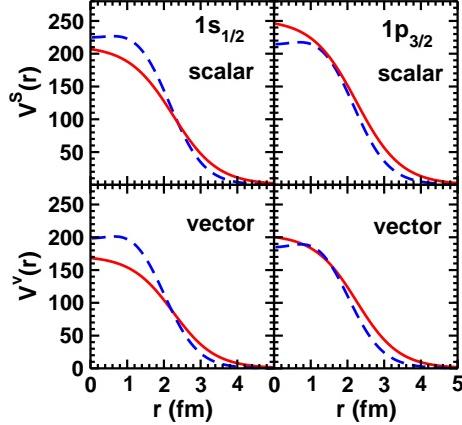


Figure 3: Scalar (upper panel) and vector (lower panel) potential fields for $1s_{1/2}$ and $1p_{3/2}$ Λ states in $^{12}\Lambda B$. The Dirac single particle and QMC model results are shown by solid and dashed lines, respectively.

where, $\rho_s(\vec{r})$ ($\rho_s^Y(\vec{r})$), $\rho_B(\vec{r})$ ($\rho_B^Y(\vec{r})$), $\rho_3(\vec{r})$ and $\rho_p(\vec{r})$ are the scalar, baryon, third component of isovector, and proton densities at the position \vec{r} in the hypernucleus [28]. On the right hand side of Eq. (8), a new, and characteristic feature of QMC appears, arrising from the internal structure of nucleon and hyperon, namely, $g_\sigma C_N(\sigma) = -[\partial M_N(\sigma)/\partial\sigma]$ and $g_\sigma^Y C_Y(\sigma) = -[\partial M_Y(\sigma)/\partial\sigma]$ where $g_\sigma \equiv g_\sigma(\sigma = 0)$ and $g_\sigma^Y \equiv g_\sigma^Y(\sigma = 0)$. The scalar and vector fields as well as the spinors for hyperons and nucleons can be obtained by solving these coupled equations self-consistently.

In Fig. 3, we compare the scalar and vector fields as calculated within the QMC model with those of the phenomenological model for $1s_{1/2}$ and $1p_{3/2}$ Λ states. It should be noted that in the QMC model the scalar and vector fields are generated by the couplings of the σ and ω mesons to the quarks. Due to the different masses of these mesons and their couplings, especially the density dependence of the σN coupling strengths, the scalar and vector fields acquire different radial dependence. In contrast to this, the two fields have the same radial shapes in the phenomenological model. We further notice that for the $1s_{1/2}$ Λ state the QMC scalar and vector fields are larger (smaller) in magnitude than those of the phenomenological model for $r < 2.5 (> 2.5)$ fm. However, for the $1p_{3/2}$ state they are smaller than the phenomenological ones everywhere.

Fig. 4 shows the moduli of the upper and lower components of $1s_{1/2}$

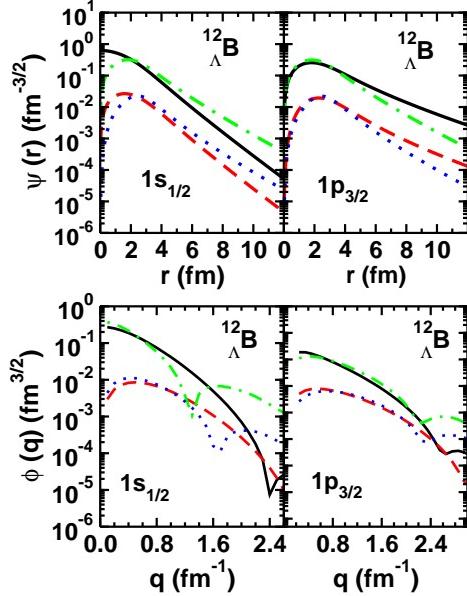


Figure 4: Moduli of upper($|f|$) and lower ($|g|$) components of the $1s_{1/2}$ and $1p_{3/2}$ Λ orbits in $^{12}_{\Lambda}\text{B}$ hypernucleus in coordinate space (upper panel) as well as in momentum space (lower panel). $|f|$ and $|g|$ of the phenomenological model are shown by solid and dashed lines, respectively while those of the QMC model by dashed-dotted and dotted lines, respectively.

and $1p_{3/2}$ Λ hyperon spinors for the $^{12}_{\Lambda}\text{B}$ in both coordinate space (upper panel) and momentum space (lower panel). We see that for the $1s_{1/2}$ Λ bound state, the spinors of the QMC model differ significantly from their phenomenological counterparts at both $r < 2$ fm and $r > 4$ fm. For the $1p_{3/2}$ Λ state while differences between them are quite big for $r > 4$ fm, this is not as prominent at smaller radii. On the other hand, in the momentum space the differences in the spinors of the two models are already quite large for $q > 1.0$ fm⁻¹ for the $1s_{1/2}$ state whereas for the $1p_{3/2}$ state the difference between the two becomes large for q beyond 2 fm⁻¹. We also note that only for $q < 1.0$ fm⁻¹, is the magnitude of the lower component ($|g(q)|$) substantially smaller than that of the upper component ($|f(q)|$). In the region of q pertinent to the kaon production, $|g(q)|$ may not be negligible. In fact, it has been shown earlier [14] that the relativistic effects resulting from the small component of Dirac bound states are large for the kaon photoproduction reactions on nuclei.

The threshold for the kaon photoproduction on ^{12}C is about 695 MeV. The momentum transfer involved in this reaction at 10° kaon angle varies between approximately 2 fm^{-1} to 1.4 fm^{-1} in the photon energy range of 0.7 GeV to 1.2 GeV [21]. In Fig. 5, we compare the differential cross section obtained by using the Λ bound state spinors calculated within the QMC and the phenomenological models for the $^{12}\text{C}(\gamma, K^+)^{12}\text{B}$ reaction. The hole state spinor was taken from the phenomenological model in both cases. The cross sections are shown for photon energies in the range of 0.7-1.2 GeV corresponding to the outgoing kaon angle of 10° . The hypernuclear states populated are 1^- , 2^- , and 2^+ , 3^+ corresponding to the particle-hole configurations of $(1p_{3/2}^{-p}, 1s_{1/2}^\Lambda)$ and $(1p_{3/2}^{-p}, 1p_{3/2}^\Lambda)$, respectively. We see that in each case the QMC cross sections are smaller than those obtained with phenomenological hyperon spinors. For the 1^- and 2^- states (involving s state Λ spinors), the QMC cross sections are lower because the corresponding momentum space spinors are smaller than their phenomenological model counterparts in the relevant momentum region. For the 2^+ and 3^+ states, additionally, the QMC potentials are also smaller than the phenomenological ones, which leads to lower QMC cross sections. In this figure we further note that the peaks of the QMC cross sections are somewhat shifted toward lower photon energies as compared to those of the phenomenological model. This can be understood from the fact that at lower photon energies the momentum transfer to the nucleus is relatively larger. In this region the QMC momentum space Λ spinors are larger as compared to those of the phenomenological model.

We further note that within each group the highest J state is most strongly excited, which is in line with the results presented in Refs. [14, 19, 46]. Furthermore, unnatural parity states within each group are preferentially excited by this reaction. The unnatural parity states are excited through the spin flip process. Thus this confirms that kaon photo- and also electro-production reactions on nuclei are ideal tools for investigating the structure of unnatural parity hypernuclear states. The addition of unnatural parity states to the spectrum of hypernuclei is expected to constrain the spin dependent part of the effective $\Lambda - N$ interaction more tightly.

In summary, we have studied the hypernuclear production by the (γ, K^+) reaction on ^{12}C within a covariant model, using hyperon bound state spinors derived from the latest quark-meson coupling model. This is the first time that quark degrees of freedom has been explicitly invoked in the description of the hypernuclear production. In our model, in the initial collision of the photon with a target proton, $N^*(1710)$, $N^*(1650)$ and $N^*(1720)$ baryonic

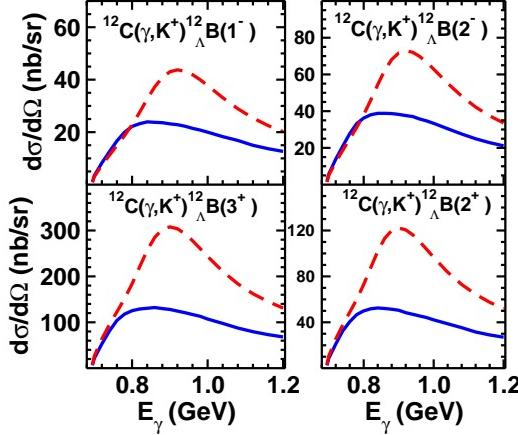


Figure 5: Differential cross sections (for the outgoing kaon angle of 10°) for the $^{12}\text{C}(p, K^+)^{12}\Lambda\text{B}$ reaction leading to hypernuclear states as indicated. The solid and dashed lines show the results of QMC and phenomenological models, respectively.

resonances are excited which subsequently propagate and decay into a Λ hyperon that gets captured in one of the nuclear orbits, while the other decay product K^+ goes out. In contrast to the previous study within this model [19], we fix the coupling constants at both the electromagnetic and hadronic resonance vertices by describing both the total and the differential cross sections of the elementary $\gamma p \rightarrow \Lambda K^+$ reaction in the relevant region of photon energies. Thus the input parameters are better constrained in this study.

We have also performed calculations with bound Λ spinors obtained by solving the Dirac equation with vector and scalar potential fields having Woods-Saxon shapes. Their depths are fitted to the binding energies of the respective states for a given set of geometry parameters which are taken to be the same for the two fields. In contrast to this model, the QMC vector and scalar fields have different radial shapes. Furthermore, both shapes and absolute magnitudes of the QMC fields are different from their Dirac counterparts. For the cases studied in this paper, the hypernuclear production cross sections calculated with the QMC hyperon spinors and fields are not only smaller in magnitude but also they peak at relatively lower photon energies as compared to those obtained within the phenomenological model.

The distortion effects in the K^+ channel have not been included in this study. However, as shown in Refs. [14, 46], these effects are weak for reactions

on p -shell nuclei but they may be more significant for heavier systems. The cross sections as calculated in this paper may be uncertain to the extent of about 10% due to the non-inclusion of the nucleon intermediate states (Born terms).

Our calculations further confirm that due to the selective excitation of the high spin unnatural parity states, the (γ, K^+) reaction on nuclei is an ideal tool for investigating the spin-flip transitions. Therefore, electromagnetic reactions provide a more complete knowledge of hypernuclear spectra and will impose more severe constraints on the poorly known spin dependent parts of the models of the $\Lambda - N$ interaction. Our model should be extended to electroproduction of hypernuclei (where the hadronic part remains the same as that discussed in this paper) so that the role of the quark degrees of freedom in the Λ bound states can be checked against the data taken at JLab.

This work has been supported by the United States Department of Energy contract no. DE-AC05-06OR23177 under which the Jefferson Science Associates (JSA) operates the Thomas Jefferson National Accelerator Facility.

Bibliography

- [1] O. Hashimoto and H. Tamura, Progr. Part. Nucl. Phys. **57** (2006) 564.
- [2] H. Bando, T. Motoba, J. Zofka, Int. J. Mod. Phys. **5** (1990) 2021.
- [3] H. Nemura, Y. Akaishi, and Y. Suzuki, Phys. Rev. Lett. **89** (2002) 142504; D. Vretenar, W. Poschl, G. A. Lalazissis, and P. Ring, Phys. Rev. C **57** (1998) 1060 (R).
- [4] C. M. Keil, F. Hoffmann, and H. Lenske, Phys. Rev. C **61** (2000) 064309; C. M. Keil and H. Lenske, Phys. Rev. C **66** (2002) 054307.
- [5] T. Miyoshi et al., Phys. Rev. Lett. **90** (2003) 232502.
- [6] L. Yuan et. al., Phys. Rev. C **73** (2006) 044607.
- [7] M. Iodice et al., Phys. Rev. Lett. **99** (2007) 052501.
- [8] F. Cusanno et al., arXiv:0810.3853 [nucl-ex].

- [9] H. Yamazaki et al., Phys. Rev. C **52** (1995) R1157; K. Maeda et al., Nucl. Phys. **A577** (1994) 277c.
- [10] J. Pochodzalla, Nucl. Phys. A754 (2005) 430c.
- [11] V. Shklyar, H. Lenske and U. Mosel, Phys. Rev. C **72** (2005) 015210.
- [12] S.S. Hsiao and S. R. Cotanch, Phys. Rev. C **28** (1983) 1668.
- [13] J. Cohen, Phys. Rev. C **37** (1988) 187.
- [14] C. Bennhold and L.E. Wright, Phys. Rev. C **39** (1989) 927; *ibid*, Phys. Lett. **B191** (1987) 11.
- [15] T. Motoba, M. Sotona, and K. Itonaga, Progr. Theo. Phys. (suppl) **117** (1994) 123.
- [16] T.-S. H. Lee, Z.-Y. Ma, B. Saghai, and H. Toki, Phys. Rev. C **58** (1998) 1551.
- [17] F.X. Lee, T. Mart, C. Bennhold, H. Haberzettl, and L. E. Wright, Nucl. Phys. **A695** (2001) 237.
- [18] R.A. Adelseck, C. Bennhold, and L.E. Wright, Phys. Rev. C **32** (1985) 1681.
- [19] R. Shyam, H. Lenske and U. Mosel, Phys. Rev. C **77** (2008) 052201(R).
- [20] R. Shyam, H. Lenske and U. Mosel, Phys. Rev. C 69 (2004) 065205; R. Shyam, H. Lenske and U. Mosel, Nucl. Phys. A764 (2006) 313.
- [21] R. Shyam, Progr. Part. Nucl. Phys. **61** (2008) 212.
- [22] P. A. M. Guichon, Phys. Lett. **B200** (1988) 235.
- [23] P. A. M. Guichon, K. Saito, E. N. Rodionov, and A. W. Thomas, Nucl. Phys. **A601** (1996) 349.
- [24] P. A. M. Guichon, H. H. Matevosyan, N. Sandulescu, A. W. Thomas, Nucl. Phys. **A772** (2006) 1.
- [25] K. Saito and A. W. Thomas, Phys. Rev. C **51** (1995) 2757.

- [26] B. D. Serot and J. D. Walecka, Adv. Nucl. Phys. **16** (1986) 1; *ibid*, Int. J. Mod. Phys. **E 6** (1997) 515.
- [27] K. Saito, K. Tsushima and A. W. Thomas, Nucl. Phys. **A609** (1996) 339.
- [28] K. Tsushima, D. H. Lu, A. W. Thomas, and K. Saito, Phys. Lett. B443 (1998) 26.
- [29] K. Tsushima, K. Saito, J. Haidenbauer and A. W. Thomas, Nucl. Phys. **A630** (1998) 691.
- [30] K. Tsushima, D. H. Lu, A. W. Thomas, K. Saito R. H. Landau, Phys. Rev. C **59** (1999) 2824.
- [31] S. D. Bass and A. W. Thomas, Phys. Lett. B634 (2006) 368.
- [32] K. Saito, K. Tsushima and A. W. Thomas, Prog. Part. Nucl. Phys. **58** (2007) 1.
- [33] J. Rikovska-Stone, P. A. M. Guichon, H. H. Matevosyan and A. W. Thomas, Nucl. Phys. **A792** (2007) 341.
- [34] P. A. M. Guichon, A. W. Thomas, and K. Tsushima, Nucl. Phys. **A 814** (2008) 66.
- [35] G. Penner and U. Mosel, Phys. Rev. C **66**, 055211 (2002); **66** (2002) 055212.
- [36] R. Shyam, Phys. Rev. C **73** (2006) 035211; R. Shyam, Phys. Rev. C **60** (1999) 055213.
- [37] T. Mart and A. Sulaksono, Phys. Rev. C **74** (2006) 055203.
- [38] R. Shyam and O. Scholten, Phys. Rev. C **78** (2008) 065201.
- [39] K.-H. Glander et al., Eur. Phys. J. **A19** (2004) 251.
- [40] R. Bradford et al., Phys. Rev. C **73** (2006) 035202.
- [41] R. A. Arndt, Ya. I. Azimov, M. V. Polyakov, I. I. Strakovsky and R. L. Workman, Phys. Rev. C **69** (2004) 035208.

- [42] T. Mart and A. Sulaksono, arXiv:nucl-th/0701007.
- [43] R. Shyam, W. Cassing and U. Mosel, Nucl. Phys. **A586** (1995) 557; R. Shyam, A. Engel, W. Cassing, and U. Mosel, Phys. Lett. **B273** (1991) 26.
- [44] A. S. Rosenthal and F. Tabakin, Phys. Rev. C **22** (1980) 711;
- [45] M. W. Ahmed et al., Phys. Rev. C **68** (2003) 064004.
- [46] A. S. Rosenthal et al., Ann. Phys. (N.Y.) **184** (1988) 33.